

Theory of Winds in Late-type Evolved and Pre-Main-Sequence Stars

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I. Introduction

A variety of recent observational results has confirmed what solar and stellar astronomers have long suspected, namely, that many of the physical processes which are known to occur in the Sun also occur among late-type stars in general. One such process is the continuous loss of mass from a star in the form of a wind. There now exists an abundance of either direct or circumstantial evidence which suggests that most (if not all) stars in the cool portion of the HR diagram possess winds. In the present brief review, an attempt is made to assess the current state of our theoretical understanding of mass loss from two distinctly different classes of late-type stars: the post-main-sequence giant/supergiant stars and the pre-main-sequence T Tauri stars. Toward this end, the observationally inferred properties of the winds associated with each of the two stellar classes under consideration are briefly summarized and then compared against the predictions of existing theoretical models. Through this analysis it will become apparent that although considerable progress has been made in attempting to identify the mechanisms responsible for mass loss from cool stars, many fundamental problems remain to be solved.

II. Mass Loss From Late-type Giants and Supergiants a. Wind Properties

The existence of winds from cool giant and supergiant stars is inferred from the detection of one or more characteristic spectral features whose formation requires that the stellar atmosphere be extended and in a state of outward expansion. Among the most frequently used such indicators are: 1) blue-shifted circumstellar absorption lines due to resonance transitions of neutral or singly-ionized metals; 2) profiles of collision-dominated chromospheric emission lines (e.g., Ca II H and K, Mg II h and k) in which the intensity of the blue K_2 (or k_2) peak is lower than that of the red K_2 (or k_2) peak; and 3), the presence of a $10\ \mu\text{m}$ emission feature and/or infrared excess, attributed to radiating silicate dust grains contained in an outflowing circumstellar gas shell. Observations of any or all of these features can in principle be analyzed to determine the mass loss rate \dot{M} of a given star. For giants and supergiants with spectral types K through M the \dot{M} values so derived are generally in the range $10^{-11} \lesssim \dot{M} \lesssim 10^{-5}\ M_{\odot}\ \text{yr}^{-1}$, with the coolest supergiant stars exhibiting the highest rates of mass loss. Unfortunately, these results are extremely model-dependent in that they are sensitive to assumptions made concerning: 1) the velocity, density, and temperature distributions throughout the wind; 2) the ionization and chemical equilibrium (or the lack thereof) in the outflowing gas; 3) the geometry of the flow and the spatial extent of the region in which a particular spectral feature is formed; and 4), the transfer of radiation in the wind. Consequently, although it can be safely said that late-type, low-gravity stars lose mass at rates significantly higher than the rate at which the Sun loses mass due to the solar wind ($\dot{M}_{\odot} \sim 10^{-14}\ M_{\odot}\ \text{yr}^{-1}$),

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quantitatively reliable estimates of \dot{M} for individual stars are presently unavailable. Indeed, because of the factors enumerated above, rates of mass loss derived for the same star by different observers can sometimes disagree by more than two orders of magnitude. A more detailed discussion of the methods and assumptions used in extracting mass loss rates from observations is contained in the review by Castor (1981).

In addition to the mass loss rate, two other observationally determined quantities are essential to the task of comparing the various theoretical models for mass loss from late-type evolved stars. These are the terminal velocity V_∞ and the temperature T of the outflow. In the case of the former, since blue-shifted absorption cores result from the scattering of photospheric radiation by atoms or ions in an expanding circumstellar envelope, estimates for V_∞ can be obtained by measuring the shortward displacements of observed lines. For those *G*, *K*, and *M* giants and supergiants whose spectra exhibit such line profiles, analyses of this type yield values for V_∞ in the range $10 \lesssim V_\infty \lesssim 100 \text{ km s}^{-1}$ (Reimers 1975; 1977), and indicate that V_∞ (on the average) decreases with decreasing stellar gravity (i.e., V_∞ is generally lowest for cool *M* supergiants). Moreover, if the measured wind terminal velocities are compared with surface gravitational escape speeds deduced from evolutionary considerations, it is found that $V_\infty \sim (0.1-0.5) V_{esc}$, where for a star of mass M_* and radius R_* , $V_{esc} = (2GM_*/R_*)^{1/2}$. This property (i.e., that $V_\infty < V_{esc}$) distinguishes the winds of cool giants and supergiants from the outflows associated with main-sequence or evolved stars located in other portions of the HR diagram. For example, the terminal velocities of the winds from luminous *O* stars follow the approximate relation $V_\infty \approx 3 V_{esc}$ (Abbott 1978), where in this case the escape speed V_{esc} includes the effect of the outward radiation force due to electron scattering opacity (see also the review by Cassinelli in this volume).

For stars having effective temperatures $T_{eff} \lesssim 6000 \text{ K}$, information concerning the temperature of the gas contained within a stellar atmosphere or wind is obtained through observations of spectral lines (either in emission or absorption) and continua formed under physical conditions similar to those found in the outer solar atmosphere. Thus, the presence of chromospheric ($T \sim 10^4 \text{ K}$) gas is inferred from the detection of lines such as Ca II H ($\lambda 3968$) and K ($\lambda 3934$), Mg II h ($\lambda 2803$) and k ($\lambda 2796$), and the $L\alpha$ ($\lambda 1216$) line of H I, while gas at transition region temperatures ($T \sim 10^5 \text{ K}$) is indicated by emission lines due to multiply-ionized species such as Si IV, C III, C IV, and N V. Likewise, evidence for the existence of coronal regions ($T \gtrsim 10^6 \text{ K}$) in the atmospheres of late-type stars is derived from attempts to observe emission at soft X-ray wavelengths. The results of numerous observational programs intended to survey the thermal properties of cool star atmospheres can be summarized as follows. While chromospheric emission lines are present in the spectra of virtually all late-type stars, transition region emission lines are not exhibited by stars whose location in the HR diagram is above and to the right of a line with approximate coordinates (K2, III), (G5 Ib) (Linksy and Haisch 1979; Simon, Linsky, and Stencel 1982). Moreover, the giant and supergiant stars which occupy this portion of the HR diagram are largely undetected as sources of coronal soft X-ray emission (Vaiana *et al.* 1981; Ayres *et al.* 1981). On this basis, we conclude that the expanding atmospheres of the coolest low-gravity stars are characterized by temperature distributions in which (for the most part) $T \lesssim 10^4 \text{ K}$. The reader is referred to the papers by Castor (1981), Dupree (1981; 1983, this volume), Linsky (1981a), and Cassinelli and Mac Gregor (1983) for additional information concerning the observationally inferred properties of mass loss from cool stars.

b. Mass Loss Mechanisms

In the present section, several of the theories which have been proposed to explain mass loss from cool giant and supergiant stars are described. The applicability of each of

the theories is measured by comparing the results of qualitative model calculations for the wind from a hypothetical K5 supergiant ($M_* = 16M_\odot$, $R_* = 400R_\odot$, $T_{\text{eff}} = 3500$ K) with the "average" wind properties (e.g., \dot{M} , V_∞ , T) derived from observations (cf. sec. IIa).

i. Thermally-Driven Winds

It is instructive to consider the possibility that the winds of late-type low-gravity stars are driven by the same mechanism which is responsible in part for the acceleration of the solar wind; namely, the force due to the thermal pressure gradient in the outflowing gas. For an isothermal, single-fluid wind model which is taken to be both spherically-symmetric and inviscid, it can be shown (see, e.g., Parker 1963; Leer, Holzer, and Flå 1982) that the flow velocity V_0 at a reference level r_0 in the stellar atmosphere is approximately

$$V_0 \approx a z_s^2 \exp\left(\frac{3}{2} - 2z_s\right) . \quad (1)$$

where for a gas of temperature T and mean mass per particle μ , $a = (kT/\mu)^{1/2}$ is the sound speed, $z_s = r_s/r_0 = (GM_*/2a^2 r_0)$ is the sonic point location (i.e., the distance r at which the wind velocity V becomes equal to a) in units of r_0 , and it has been assumed that $V_0 \ll a$. From conservation of mass, it follows that the rate of mass loss is given by $\dot{M} = 4\pi r_0^2 \mu N_0 V_0$ where N_0 is the total number density at r_0 . To apply equation (1) to a star with the physical parameters given above, note that the absence of detectable transition region and coronal emission implies $T \lesssim 10^4$ K (cf. sec. IIa). Adopting $T = 10^4$ K, $r_0 = R_*$, and $\mu = 0.667 m_H$ (the value appropriate to a plasma composed of H and He with $N_e = N_H$, $N_{He}/N_H = 0.1$), it is found that $z_s \approx 31$ and $V_0 \approx 8 \times 10^{-18}$ cm s $^{-1}$, indicating that for physically reasonable values of N_0 mass loss at rates $\gtrsim 10^{-11} M_\odot$ yr $^{-1}$ cannot be thermally-driven. This conclusion is not particularly dependent upon the assumed isothermal temperature distribution (cf. Leer, Holzer, and Flå 1982), and can also be reached in the following equivalent way (Weymann 1962; 1978). For a given value of N_0 we ask, what is the value of T such that the initial velocity V_0 calculated from equation (1) yields a particular mass loss rate \dot{M} ? Assuming $N_0 = 10^{11}$ cm $^{-3}$, the temperature required to produce a wind having $\dot{M} = 10^{-8} M_\odot$ yr $^{-1}$ from a star with the hypothetical K5 supergiant parameters is $T \approx 70,000$ K. Such a temperature is incompatible with the observed lack of transition region emission lines (e.g., these due to C IV) in the spectra of cool supergiants, since for a gas in collisional ionization equilibrium at $T \approx 70,000$ K, approximately 10% of all C is in the form of C IV (Jordan 1969).

ii. Radiatively-Driven Winds

Numerous authors have suggested that the winds of cool giants and supergiants can be driven by the force arising from the scattering and absorption of radiation from the stellar photosphere by sources of opacity in the outflow. Although this mechanism appears to be responsible for mass loss from luminous early-type stars (see e.g., the review by Cassinelli in this volume), some fundamental problems are encountered in attempting to account for the outflows from late-type stars in terms of a radiation-driven wind model. One such difficulty arises from the fact that the atoms and ions known to be prevalent in the winds of cool stars have strong resonance lines located in the visible and UV portions of the spectrum, while the photospheric continuum radiation field is most intense in the red or near IR (Goldberg 1979). Several efforts have been made to circumvent this spectral mismatch, including the use of molecular opacity (Maciel 1976, 1977) and the force due to the scattering of chromospheric $L\alpha$ radiation by hydrogen atoms in the flow (Wilson 1960; Haisch, Linsky, and Basri 1980). Both of these models are unsatisfactory in that the small values of the respective opacities requires inordinately large stellar luminosities to produce outflows with momentum fluxes $\dot{M} V_\infty$ comparable to those inferred from observations (see the discussion of Mac Gregor 1982). In the remainder of

this section we consider a third proposal, namely, that the condensation of dust grains and their subsequent outward acceleration by radiative forces can cause the entire circumstellar envelope to expand into a wind (Gehrz and Woolf 1971).

To investigate the efficacy of this mechanism, consider a spherically-symmetric wind model in which the collisional coupling of the grains to the background gas is sufficiently strong (see Gilman 1972) to ensure that the optically-thin radiative force per unit volume exerted on the envelope can be written as

$$f_{\text{rad}} = \frac{K_d \rho_g L_*}{4\pi c r^2} \quad (2)$$

where L_* is the stellar luminosity, ρ_g is the gas mass density, and

$$K_d = \frac{3Q_{\text{pr}}}{4r_{\text{gr}}\rho_{\text{gr}}} \cdot \frac{\rho_d}{\rho_g} \quad (3)$$

is the dust opacity. In equation (3), ρ_d is the dust mass density, r_{gr} and ρ_{gr} are, respectively, the radius and density of an individual grain, and Q_{pr} is the radiation pressure efficiency factor. For illustrative purposes, adopt $\rho_d/\rho_g = 10^{-2}$ and consider grains of radius $r_{\text{gr}} = 10^{-5}$ cm composed of the common silicate olivine (Mg_2SiO_4) for which $\rho_{\text{gr}} = 3.2 \text{ g cm}^{-3}$. For pure ("clean") silicate grains, $Q_{\text{pr}} \sim 0.1$ (Gilman 1974) and $K_d \approx 23.4 \text{ cm}^2 \text{ g}^{-1}$, while for silicate grains containing impurities (i.e., "dirty" grains) Q_{pr} may be as much as a factor of 10 larger (using the results of Gilman [1974] for graphite) yielding $K_d \approx 234.0 \text{ cm}^2 \text{ g}^{-1}$. To simulate the formation of grains in the flow, assume that $K_d = 0$ at a reference level r_0 ($\approx R_*$) and increases linearly in magnitude to either of the values given above within a distance Δr , remaining constant for $r > r_0 + \Delta r$. Under these conditions, the equation of motion for an isothermal wind can be integrated to obtain the following approximate expressions for the critical point location z_c ($= r_c/r_0$), initial velocity V_0 , and terminal velocity V_∞ :

$$z_c = z_s \frac{1 + \Gamma/\Delta z}{1 + \Gamma z_s/\Delta z} \quad (4)$$

$$V_0 \approx a \exp(-z_s \Delta z / \Gamma) \quad (5)$$

$$V_\infty \approx V_{\text{esc}}(\Gamma - 1)^{1/2} \quad (6)$$

In equations (4) through (6), a is the sound speed, V_{esc} is the escape speed at r_0 , z_s is the sonic point location, and $\Gamma = (K_d L_* / 4\pi G M_* c)$ is the ratio of the radiative force to the gravitational force acting on the gas. To apply these results, consider a star having the K5 supergiant parameters given above, and assume $r_0 = R_*$, $\Delta z = 0.1$, $N_0 = 10^{11} \text{ cm}^{-3}$, $T = 3000 \text{ K}$ ($\approx 0.85 \text{ Teff}$). In the case of "clean" silicate grains $\Gamma = 2.43$ and equations (4) through (6) yield $z_c = 1.041$, $\dot{M} \approx 4.5 \times 10^{-9} M_\odot \text{ yr}^{-1}$, $V_\infty \approx 1.2 V_{\text{esc}}$ ($= 148 \text{ km s}^{-1}$), while for "dirty" grains $\Gamma = 24.3$ with $z_c = 1.004$, $\dot{M} \approx 6.5 \times 10^{-8} M_\odot \text{ yr}^{-1}$, and $V_\infty \approx 4.8 V_{\text{esc}}$ ($= 596 \text{ km s}^{-1}$). These qualitative results indicate that values of \dot{M} in accord with observations can be produced by the mechanism, a conclusion which is substantiated by more detailed calculations (Salpeter 1974; Kwok 1975; Goldreich and Scoville 1976; Lucy 1976; Menietti and Fix 1978; Philips 1979). Moreover, the computed values of V_∞ may be reduced through the inclusion of grain destruction through sputtering (cf. Kwok 1975).

At this point, it is important to consider whether or not the schematic picture of grain formation adopted above is consistent with the derived wind models. To investigate this question, we follow the analysis of Draine (1981) who has noted that condensation

requires (among other things) the vibrational temperature T_v of small olivine clusters to be less than a saturation value $T_{v,sat.}$ given implicitly by the relation

$$T_{v,sat.} = \frac{6.13 \times 10^4 K}{\ln \left[\frac{1.16 \times 10^{35} \text{ cm}^{-3}}{N} \left(\frac{T_{v,sat.}}{T} \right)^{\frac{3}{2}} \right]} \quad (7)$$

Solution of equation (7) for physical conditions appropriate to the critical point in the "clean" grain model ($N_c \approx 2.8 \times 10^7 \text{ cm}^{-3}$, $T = 3000 \text{ K}$) given above yields $T_{v,sat.} = 973 \text{ K}$. The actual temperature T_v of silicate grains with radius $r_{gr} = 10^{-5} \text{ cm}$ at the critical point can be straightforwardly determined from a simple description of the grain energy balance which includes the effects of radiative and collisional heating and cooling (cf. Draine 1981, eq. [22]). For the derived critical point conditions T_v is controlled by radiative processes and has the value $T_v = 2475 \text{ K} > T_{v,sat.}$, indicating that grain formation *cannot* occur in the manner assumed.

Since the grain energy balance is dominated by radiative heating and cooling, condensation is possible (for sufficiently high densities) only at larger distances where the stellar radiation field is dilute. Draine (1981) has shown that the nucleation of small ($r_{gr} \sim 3 \text{ \AA}$), "clean" silicate grains can take place within several stellar radii of the photosphere when $T_{eff} \leq 3500 \text{ K}$, but for grains of this size, $Q_{pr} < 10^{-5}$ (Gilman 1974) making them dynamically unimportant. Hence, because formation of grains with desirable optical properties cannot occur close enough to the stellar surface to ensure production of a wind with a mass loss rate $\gtrsim 10^{-9} M_\odot \text{ yr}^{-1}$, we conclude that the mechanism is unable (by itself) to account for the winds from the majority of cool, low-gravity stars. The reader is referred to the reviews by Castor (1981), Linsky (1981b), Mac Gregor (1982), and Casinelli and Mac Gregor (1983) for additional discussion of dust-driven wind models.

iii. Shock Wave-Driven Winds

Observational evidence for mass loss from long period (Mira) variables has led several investigators (Willson 1976; Slutz 1976; Wood 1979; Willson and Hill 1979) to suggest that the outflows associated with these stars may be produced by the pulsations responsible for their brightness variations. The physical mechanism by means of which mass is ejected from the stellar atmosphere can be understood as follows. The periodic oscillations of the surface of the star give rise to compressional disturbances which propagate outward in the form of shock waves. In the lower atmospheric layers where the gas density is high and radiative cooling is efficient, these travelling shocks behave isothermally. Consequently, each parcel of shocked gas tends to return to its pre-shock dynamical state following the passage of discontinuity. However, higher in the atmosphere where the gas density is lower, the shocks become adiabatic and the energy gained by the post-shock gas can be sufficient to produce mass loss.

The discussion given above suggests that an estimate of the mass loss rate due to a shock wave-driven wind can be obtained if the density at the atmospheric level above which the shocks behave adiabatically is known (Willson and Hill 1979; Castor 1981). Assuming that the volume radiative cooling rate of the post-shock gas can be written as $N^2 P_R(T)$ where $P_R(T)$ is the cooling coefficient, the cooling time is of order $t_{cool} \sim kT / NP_R(T)$. Similarly, if the flow time for the post-shock gas is taken to be the time required to travel a distance equal to the local density scale height, then $t_{flow} \sim a/g$ where $g = GM_\star / r^2$ is the gravitational acceleration and it has been assumed that the flow velocity is the sound speed a . The transition from isothermal to adiabatic shock behavior occurs at the level where $t_{cool} \sim t_{flow}$; equating derived expressions for the flow and cooling times yields $N \sim gkT / aP_R(T)$, from which it follows that the maximal mass loss rate is

$$\dot{M} \lesssim 4\pi r^2 \mu Na \sim 4\pi GM_* \mu kT / P_R(T) \quad (8)$$

According to equation (8), \dot{M} is inversely proportional to the radiative cooling coefficient $P_R(T)$, a quantity which should be determined from a consistent transfer calculation. To obtain an order of magnitude estimate for \dot{M} , however, we assume that the emitting post-shock material is characterized by a chromospheric gas temperature (cf. Wood 1979; Willson and Hill 1979; Willson and Pierce 1982) and evaluate $P_R(T)$ from the approximate chromospheric cooling law given by Hartmann, MacGregor, and Avrett (1983). Using this prescription for $T = 10^4$ K, $P_R(T) \approx 2.6 \times 10^{-24}$ erg cm³ s⁻¹ and $\dot{M} \lesssim 1.6 \times 10^{-11} (M_*/M_\odot) M_\odot \text{yr}^{-1}$. If the post-shock gas temperature is as low as $T = 5000$ K, then $P_R(T) \approx 2.2 \times 10^{-26}$ erg cm³ s⁻¹ and $\dot{M} \lesssim 1.8 \times 10^{-9} (M_*/M_\odot) M_\odot \text{yr}^{-1}$. The extreme sensitivity of these results demonstrates the need for detailed calculations including transfer effects in order to determine whether or not substantial mass loss can be driven by this mechanism; preliminary steps in this direction have been taken by Willson and Pierce (1982). We further note that although the shock wave-driven wind model has been developed primarily to explain mass loss from Mira variables, virtually all M supergiants are semi-regular or irregular variables (see, e.g., Feast 1981). However, it is not known at the present time whether this variability is the result of actual pulsation or is due instead to the motions of large convective elements on the stellar surface (Schwarzschild 1975). Moreover, as Linsky (1981b) has noted, the mechanism is probably unable to account for the winds from stars having lower luminosities and higher surface gravities than the Mira variables (e.g., the G-K supergiants) since there is no observational evidence for the presence of large amplitude shock waves in the atmospheres of these stars.

iv. Alfvén Wave-Driven Winds

The effects of an outwardly propagating flux of Alfvén waves on the dynamics of winds from cool giants and supergiants has been considered by Hartmann and MacGregor (1980, 1982a). This investigation was motivated by the fact that direct observations of the solar wind plasma and the interplanetary magnetic field typically reveal the presence of hydromagnetic fluctuations, many of which appear to be Alfvénic in character (see, e.g., the review by Barnes 1979). Numerous authors have suggested that such wave modes may be responsible for providing the additional energy required to produce high-speed streams in the solar wind (Belcher 1971; Alazraki and Couturier 1971; Hollweg 1973, 1978; Belcher and Olbert 1975; Jacques 1977, 1978; Leer, Holzer, and Flå 1982).

To examine the properties of a wind driven by Alfvén waves, consider a steady, spherically-symmetric outflow which emanates from a star having a radially-directed magnetic field B . Assume that the wave amplitude δB is everywhere smaller than B , and that the wavelength of the fluctuation is shorter than any of the scale lengths over which the wind properties vary (i.e., the WKB approximation). The frequency ω of a wave which propagates in the radial direction is then a constant and is given in terms of the wave vector k , wind velocity V , and Alfvén speed $A = B(4\pi\rho)^{-1/2}$ by the dispersion relation $\omega = k(V + A)$. For a wave of energy density $\varepsilon = \delta B^2/8\pi$ which propagates without attenuation, a Lagrangian treatment of the wave properties (see, e.g., Jacques 1977) indicates that the action density $S = \varepsilon/(\omega - kV)$ is conserved in the sense that $\nabla \cdot (\mathbf{V}_g S) = 0$, where $\mathbf{V}_g = \mathbf{V} + \mathbf{A}$ is the group velocity. From the definition of S and the assumed spherical symmetry of the outflow, it follows that $\varepsilon \propto [M_A(1 + M_A^2)]^{-1}$, where $M_A = V/A \propto \rho^{-1/2}$ is the Alfvénic Mach number. The time-averaged force exerted by the wave on the moving background plasma through which it travels is simply $-\frac{1}{2} d\varepsilon/dr$; since ε decreases with distance from the star, the wave exerts an outward force on the gas. Physically, the local reduction in ε is accompanied by an increase in the streaming energy per unit mass of the wind. This can be seen by straightforwardly calculating the divergence of the wave

energy flux $E_A = \varepsilon (\frac{3}{2}V + A)$ (cf. Belcher 1971), and recognizing that the non-zero result is equal to the rate at which the wave does work on the flow.

An approximate solution to the equation of motion for a wave-driven wind can be obtained for the limit in which the wave force dominates the thermal pressure gradient force (cf. Leer, Holzer, and Flå 1982). Defining $\beta = \varepsilon_0 / (\frac{1}{2} \rho_0 V_{esc}^2)$ (the subscript "0" denotes evaluation at a reference level r_0), the resulting expressions for the critical point location r_c , initial velocity V_0 , and terminal velocity V_∞ are

$$r_c / r_0 \approx \frac{7}{4(1 + \beta/2)} . \quad (9)$$

$$V_0 / V_{esc} \approx \frac{1}{8} \beta^2 \left[\frac{7}{4(1 + \beta/2)} \right]^{7/2} . \quad (10)$$

$$V_\infty / V_{esc} \approx (\beta / M_{A0} - 1)^{1/2} . \quad (11)$$

where equation (11) is applicable for β , $V_0 / V_{esc} \ll 1$. For illustration, adopt the K5 supergiant parameters given earlier in this section and assume $r_0 = R_*$, $N_0 = 10^{11} \text{ cm}^{-3}$, $\mu = 0.867 m_H$, $B_0 = 10 G$, $\delta B_0 = B_0 / \sqrt{10}$ (corresponding to $\beta = 4.7 \times 10^{-2}$, $F_{A0} = 3.36 \times 10^6 \text{ erg cm}^{-2} \text{ s}^{-1}$). For this specific (but arbitrary) choice of reference level location and parameters, equations (10) and (11) then yield $\dot{M} = 4\pi r_0^2 \mu N_0 V_0 \approx 3.8 \times 10^{-7} M_\odot \text{ yr}^{-1}$ and $V_\infty \approx 4.1 V_{esc} = 509 \text{ km s}^{-1}$. These qualitative results are representative of those obtained from more detailed calculations (cf. Hartmann and Mac Gregor 1980; Holzer, Flå, and Leer 1983): mass loss rates in the range $10^{-9} - 10^{-6} M_\odot \text{ yr}^{-1}$ can be obtained for $F_{A0} \sim 10^5 - 10^6 \text{ erg cm}^{-2} \text{ s}^{-1}$, but because the waves increase the streaming energy of the gas in the supersonic portion of the outflow, they give rise to winds having $V_\infty > V_{esc}$, in contradiction to observations.

It was noted by Hartmann and Mac Gregor (1980) that the terminal velocity V_∞ of an Alfvén wave-driven wind could be reduced if it was assumed that the waves were dissipated at a rate such that the e-folding length L for the decrease in wave energy density (the damping length) was constant and equal in magnitude to the stellar radius. Such a prescription causes most of the wave energy to be deposited near the base of the flow, resulting in winds with mass loss rates which are nearly the same as those obtained for undamped waves, but which have terminal velocities $V_\infty < V_{esc}$. An important consequence of the required wave dissipation is that the flow is heated at a rate $\Gamma = \varepsilon(V + A)/L$, leading to the production of a region containing gas at chromospheric temperatures which extends several stellar radii above the base of the wind. That such extended chromospheres do in fact exist about late-type giants and supergiants is indicated by a variety of recent observations (Stencel 1982 and reference therein; Newell and Hjellming 1982; Hjellming and Newell 1983; Beckers *et al.*, 1983). Unfortunately, recent more detailed studies indicate that the conditions under which Alfvén wave dissipation can be described in terms of a constant damping length are unlikely to be realized in the winds of cool, low-gravity stars. Holzer, Flå, and Leer (1983) have computed wind models which include a consistent treatment of the frictional wave damping due to ion-neutral collisions in the partially-ionized gas. They find that the wave-induced changes in the background flow properties cause the frictional damping length to vary in size by several orders of magnitude, resulting in (among other things) winds for which $V_\infty > V_{esc}$. These authors have also shown (see also Leer, Holzer, and Flå 1982) that even if L is taken to be constant, winds with $V_\infty \sim (0.1-0.5)V_{esc}$ correspond to an unrealistically narrow range of assumed L values. Hence, it would appear that the model as it presently stands suffers from a fundamental (and not easily correctable) deficiency in that it is unable to produce outflows having terminal velocities which are a small fraction of V_{esc} .

v. Magnetic Reconnection-Driven Winds

A final mass loss mechanism involves physical processes which are thought to occur when magnetic flux (in the form of bipolar magnetic loops) emerges through the surface of the Sun. This model has been applied to the acceleration of the solar wind (Pneuman 1983) and the winds of cool giant and supergiant stars (Mullan 1980, 1981, 1982), and bears a resemblance to the so-called "melon seed" mechanism for the ejection of flare-associated surges (Schlüter 1957; Svestka 1976). It can be described in the following way. Distortion of an ambient magnetic field by the upward motion of an emerging flux loop gives rise to restoring forces which cause the footpoints of the loop to be pinched off. Rapid, small-scale reconnection leads to the formation of a self-contained diamagnetic "plasmoid". If the plasma β is small compared to unity and if the strength of the external field decreases with height, then this element experiences an outward magnetic force, similar to the force experienced by a localized current distribution when placed in a non-uniform external magnetic field (Parker 1957). For a geometrically thin plasmoid containing gas of temperature T and density ρ which is assumed to be in pressure equilibrium with the surrounding external field B (i.e., $B^2/8\pi = 2kT\rho/m_p$), the magnitude of this force is (Pneuman 1983)

$$F = - \frac{3M}{2\rho} \frac{d}{dr} \left[\frac{B^2}{8\pi} \right] \quad (12)$$

where M is the total mass of the gas in the plasmoid.

It is of interest to determine the strength B of the external magnetic field required in order that the force given by equation (12) exceed the gravitational force acting on a plasmoid. To do this, we follow Pneuman (1983) and assume that $B \propto r^{-n}$, so that $F = (3nMB^2/8\pi\rho r)$. It then follows that $F > GM_\bullet M/r^2$ at $r = R_\bullet$ when $B > (8\pi GM_\bullet \rho/3nR_\bullet)^{1/2}$. Adopting $n = 2$ and $N = 10^{10} \text{ cm}^{-3}$ along with the physical parameters of the previously considered K5 supergiant, this criterion yields $B > 2.3G$. Although this value is not unreasonable, note that the plasmoid gas temperature required by the condition of pressure equilibrium is $T = (B^2/16\pi Nk) \approx 77,000 \text{ K}$. In view of the absence of transition region emission lines from the spectra of late-type supergiants (cf. sec. IIa), this temperature is clearly too high if most of the observed mass loss is in the form of such discrete plasmoids.

III. Mass Loss From Pre-Main-Sequence (T Tauri) Stars a. Physical Properties of T Tauri Stars

The group of objects known as T Tauri stars are thought to be the progenitors of stars like the sun, still in the process of contracting to the main sequence. They are probably among the youngest stellar objects, and have been defined by Herbig (1962) as stars whose spectra contain the Balmer lines of hydrogen (in emission and/or absorption), and emission lines due to Ca II, Fe I, and Fe II. They are variable, often exhibit ultraviolet and/or infrared excesses, and are frequently associated with low-density nebulosities. On the basis of optical and infrared observations of several hundred such objects, Cohen and Kuhi (1979) have concluded that most T Tauri stars have spectral types from mid to late K, and luminosity classes between III and V. Comparison with calculated pre-main-sequence evolutionary tracks suggests that the masses and radii of these stars are typically in the ranges $M^* \sim (0.2-3)M_\odot$ and $R_* \sim (1-5)R_\odot$, respectively. Their average age is $\sim 10^6$ years, a value comparable to the initial free-fall time for the spherical collapse of a $1M_\odot$ protostellar cloud (Larson 1978). The luminosities of T Tauri stars are generally $\lesssim 5L_\odot$, with an average effective temperature $T_{\text{eff}} \sim 4000 \text{ K}$.

b. Wind Properties

The primary evidence for mass loss from T Tauri stars comes from the detection of spectral lines having P Cygni-type profiles. These profiles consist of a blue-shifted absorption feature together with a redward-displaced emission peak, and are interpreted as arising from the scattering of stellar radiation by a spectral line formed in an expanding atmosphere (see, e.g., the review by Cassinelli in this volume for a discussion of line formation in a moving medium). From an analysis of spectra obtained by Herbig (1977) for approximately 75 T Tauri stars, Kuhl (1978) has drawn the following conclusions concerning the frequency of occurrence of different types of $H\alpha$ line profiles. About 5% of the stars observed exhibit the classical type I P Cygni line profiles in which the bottom of the blue-displaced absorption feature is well below the level of the nearby continuum; such profiles constitute unambiguous evidence for the presence of an outflow. Another 10% of the stars studied have single $H\alpha$ emission peaks which are symmetric about the laboratory wavelength of the line. Of the remaining stars, most (~60% of the total) exhibit $H\alpha$ line profiles of the P Cygni type III, in which the blue-shifted absorption does not go below the continuum level and is bounded on the short wavelength edge by a second emission peak of lower intensity than the one situated to the red of the rest wavelength of the line. The displacement of this blue emission feature generally corresponds to velocities in the range 100-200 km s⁻¹ toward the observer along the line of sight to the star, with the red emission peak having velocities 50-150 km s⁻¹ in the opposite direction (Kuhl 1978).

Despite the fact that there is considerable uncertainty regarding the way in which spectral lines having type III P Cygni profiles are formed, the presence of a blue-shifted absorption dip has led many investigators to adopt them as a signature of mass outflow. For example, Kuhl (1964, 1966) has analyzed the $H\alpha$ line profiles in the spectra of 8 T Tauri stars using a model in which the observed emission arises from material which is ballistically ejected (and subsequently gravitationally decelerated) from the stellar surface. This emitting region is surrounded by a layer of cool, absorbing gas which expands radially with constant velocity in order to produce a shortward-displaced absorption feature. By attempting to fit the observed line profiles, Kuhl derived mass loss rates in the range $2.5 \times 10^{-8} \lesssim \dot{M} \lesssim 5.8 \times 10^{-7} M_{\odot} \text{yr}^{-1}$ for the stars in his sample, with an average mass loss rate of $\dot{M} \sim 3 \times 10^{-8} M_{\odot} \text{yr}^{-1}$. Given the schematic character of the model used to interpret the data, it is difficult to assess the accuracy of Kuhl's mass loss rate estimates. However, De Campli (1981) has noted that a determination of \dot{M} from the analysis of a single feature (e.g., an asymmetric $H\alpha$ emission profile) in the spectrum of a T Tauri star is subject to the same uncertainties enumerated earlier in connection with the inferred mass loss rates of cool giants and supergiants (cf. sec. IIa). Because of this, the values of \dot{M} deduced from observed emission line intensities can differ by as much as three orders of magnitude, with the rate of mass loss from an "average" T Tauri star only constrained to lie somewhere within the range $3 \times 10^{-9} \lesssim \dot{M} \lesssim 10^{-6} M_{\odot} \text{yr}^{-1}$ (De Campli 1981).

An alternative interpretation of the type III P Cygni profiles seen in the spectra of T Tauri stars has been proposed by Ulrich (1976). He has noted that the collapse of an interstellar gas cloud to form a star is decidedly non-homologous, with the central portions of the cloud contracting rapidly to make a largely hydrostatic core, onto which a tenuous outer envelope slowly accretes (see, e.g., Larson 1969). If both the core and envelope rotate, the infalling gas must pass through an oblique accretion shock located near the surface of the core. Material traversing this discontinuity is compressed and heated, thereby producing the observed Balmer line emission. Moreover, because matter which impacts the far hemisphere of the core moves toward the observer while that striking the near hemisphere moves away, blue-shifted and red-shifted emission features are formed. The blue-shifted absorption dip between the red and blue emission peaks is a

consequence of the fact that for some velocities, radiating post-shock gas is not visible to an outside observer (cf. Ulrich 1976). As it stands, the model utilizes an approximate treatment of the transfer of radiation within the emitting region, and neglects entirely the effects of the outer portions of the infalling envelope and an equatorial accretion disk on the observed radiation (Ulrich 1978). However, there does exist observational evidence which suggests that inflows can occur in association with some T Tauri stars. For example, the higher Balmer series lines in the spectra of those T Tauri stars typified by the star YY Orionis are observed to have *inverse* P Cygni profiles (i.e., containing a redward-displaced absorption component) (Walker 1972). Estimates for the fraction of T Tauri stars which are of the YY Ori type range between 5 and 50% (Kuhi 1978; Appenzeller 1978), although no inverse P Cygni profiles were seen in the spectra of 50 T Tauri stars observed by Herbig (1977), including the star YY Ori itself. This result (and others like it; cf. Kuhi 1978; Hartmann 1982) suggests that the physical conditions which characterize the atmospheres of T Tauri stars are highly variable. To further complicate matters, the envelopes would also appear to be quite inhomogeneous since some T Tauri stars exhibit Balmer line profiles signifying the presence of outflowing gas *and* Na D line profiles indicative of infalling material (Ulrich and Knapp 1979).

Few conclusive results are available regarding the terminal velocities and temperatures of the outflows from T Tauri stars. For those stars which exhibit fairly unambiguous evidence for mass loss (e.g., T Tau itself), measurements of the blueshifts of observed absorption and/or emission features generally yield velocities $V \sim \text{several} \times 10^2 \text{ km s}^{-1}$. If it is assumed that the velocities so derived represent the terminal velocity V_∞ of a wind, then for the stellar masses and radii inferred from the HR diagrams constructed by Cohen and Kuhi (1979; cf. sec. IIIa) $V_\infty \sim V_{esc} = (2GM_\star/R_\star)^{1/2}$. Evidence for the existence of atmospheric regions in which $T > T_{eff}$ is obtained from the detection of numerous emission lines (cf. sec. IIa) due to species such as H I, He I, Ca II, Mg II, Si IV, C III, C IV, and N V in ultraviolet, visible, and infrared spectra of some well-studied T Tauri stars (see, e.g., Gahm *et al.* 1979; Cram, Giampapa, and Imhoff 1980; Imhoff and Giampapa 1980; Herbig and Soderblom 1980; Ulrich and Wood 1981; Giampapa *et al.* 1981). Such emission lines are indicative of gas at chromospheric ($T \sim 10^4 K$) and transition region ($T \sim 10^5 K$) temperatures, and in the case of T Tauri stars, measured surface fluxes in individual lines are frequently more than 10 to 100 times the corresponding average solar values. Flux enhancements of this magnitude are more characteristic of solar active regions, and may be due in part to the onset of the low chromospheric temperature rise at a much deeper atmospheric level than in the Sun (Herbig 1970; Cram 1979). Furthermore, approximately one-third of the T Tauri stars observed with The Imaging Proportional Counter on the *Einstein* satellite have been detected as soft X-ray sources, with luminosities in the range $L_X \sim 10^{30} - 10^{31} \text{ erg s}^{-1}$ (Gahm 1980; Feigelson and DeCampli 1981). It has been proposed that this emission arises from coronal ($T \gtrsim 10^6 K$) gas occupying a region near the photosphere whose spatial extent is small in comparison with the dimensions of the extended circumstellar envelope formed by the wind (Gahm 1980; Feigelson and DeCampli 1981; Walter and Kuhi 1981; Kuhi 1982). In this picture, the absence of detectable soft X-ray emission from all T Tauri stars may be the result of X-ray attenuation by the cooler, H α -emitting wind material. However, Montmerle *et al.* (1983) have used the *Einstein* Observatory to study the ρ Oph dark cloud, a region in which star formation is believed to be occurring. Repeated observations of portions of the cloud on a variety of time scales have revealed the presence of approximately 50 highly-variable soft X-ray sources, most of which are probably pre-main-sequence objects. Hence, the failure of earlier surveys to detect more T Tauri stars as X-ray sources may be a consequence of the fact that the high-temperature emission from these objects is dominated by strong, flare-like events and is discernible only from longer or repeated observations.

We conclude this section by noting that while asymmetric or blue-shifted spectral line profiles provide fairly direct information regarding the dynamical state of T Tauri

star atmospheres, there is additional indirect evidence for mass loss from these stars (see, e.g., Cohen 1982). For example, the proper motions and optical emission of Herbig-Haro objects (nebulous, "semi-stellar" knots frequently associated with T Tauri stars) have been attributed to the interaction between a condensation and the wind from a nearby pre-main-sequence star (Schwartz and Dopita 1980; Herbig and Jones 1981; Dopita, Schwartz, and Evans 1982). Likewise, recent observations of T Tauri stars at radio wavelengths have resulted in several detections of continuum (free-free) emission (Cohen, Bieging, and Schwartz 1982) and high-velocity molecular (CO) gas (Edwards and Snell 1982), both presumably arising from an extended, outflowing circumstellar envelope.

c. Mass Loss Mechanisms

As was done in the case of cool giants and supergiants, the predictions of theoretical models for mass loss from T Tauri stars are now compared against the observationally inferred properties described in the preceding section. For this purpose, we consider the wind from a hypothetical T Tauri star having physical parameters $\dot{M}_* = 0.75 M_\odot$, $R_* = 4 R_\odot$, and $T_{eff} = 4000\text{K}$.

i. Thermally-Driven Winds

In order to ascertain whether or not T Tauri stars can undergo significant thermally-driven mass loss, the isothermal, spherically-symmetric wind model described in section II b is adopted (see also Bisnovaty-Kogan and Lamzin 1977; Ulrich 1978; DeCampli 1981). We first note that according to the results of that section, mass loss rates in the range $3 \times 10^{-9} \lesssim \dot{M} \lesssim 10^{-8} M_\odot \text{ yr}^{-1}$ (DeCampli 1981; cf. sec. III b) cannot be obtained if $T \lesssim 3 \times 10^4 \text{ K}$ throughout the wind as the observed Balmer line emission would appear to indicate. However, this difficulty can be alleviated if it is assumed that the subsonic portion of the outflow is characterized by coronal gas temperatures, as suggested by the fact that some T Tauri stars have been detected as soft X-ray sources (cf. sec. III b). Assuming $T = 10^6 \text{ K}$, $\mu = 0.609 m_H$ (the value appropriate to a fully-ionized plasma composed of H and He) and using the T Tauri star parameters given at the beginning of this section, the sonic point location (in units of the stellar radius) is $z_s = (GM_* \mu / 2kT R_*) \approx 1.32$. From equation (1), it then follows that the velocity V_0 at a reference level $r_0 (= R_*)$ in the stellar atmosphere is $V_0 \approx 0.56a$, yielding $\dot{M} = 4\pi R_*^2 \mu N_0 V_0 \approx 1.02 \times 10^{-7} M_\odot \text{ yr}^{-1}$ for $N_0 = 10^{12} \text{ cm}^{-3}$. While this value for \dot{M} falls within the range dictated by observations, the model can be ruled out on the basis of the heating rate required to maintain a coronal region having the stipulated temperature and density. To see this, assume for simplicity that the corona is hydrostatic and isothermal, and occupies a volume $4\pi R_*^2 h$, where $h = (kTR_*^2 / GM_* \mu)$ is the density scale height. For an optically-thin gas in collisional ionization equilibrium, the radiative cooling rate Λ is expressible in the form $\Lambda = N_e N_H P_R(T)$ (units: $\text{erg cm}^{-3} \text{ s}^{-1}$), from which, it follows that the coronal luminosity is $L_c \approx 4\pi R_*^2 h N_e N_H P_R(T)$. Adopting $P_R(T) \approx 10^{-22} \text{ erg cm}^3 \text{ s}^{-1}$ (Raymond, Cox, and Smith 1976) for $T = 10^6 \text{ K}$, it is readily seen that $L_c \approx 2.4 \times 10^{36} \text{ erg s}^{-1}$, implying a rate of energy addition to the corona far in excess of the stellar luminosity $L_* \approx 1.4 \times 10^{34} \text{ erg s}^{-1}$. Moreover, for the assumed coronal temperature a significant fraction of L_c is radiated in the form of soft X-rays, in contradiction to the observational result that the X-ray luminosities of T Tauri stars are generally in the range $L_x \sim 10^{30} - 10^{31} \text{ erg s}^{-1}$ (cf. sec. III b). Note that this problem cannot be ameliorated by reducing T (and thereby, \dot{M}) to soften the spectrum of emitted coronal radiation. Since $P_R(T)$ increases with decreasing T for $10^5 \lesssim T \lesssim 10^6 \text{ K}$, the power input required to maintain a cooler corona is actually comparable to the value derived above. For example, adopting $T = 350,000 \text{ K}$ with $N_0 = 10^{12} \text{ cm}^{-3}$, the preceding analysis yields $z_s \approx 3.77$, $V_0 \approx 0.034a$, and $\dot{M} \approx 3.67 \times 10^{-9} M_\odot \text{ yr}^{-1}$. However, the cooling coefficient corresponding to this temperature is $P_R(T) \approx 3.25 \times 10^{-22} \text{ erg cm}^3 \text{ s}^{-1}$, so that $L_c \approx 2.7 \times 10^{36} \text{ erg s}^{-1}$.

ii. Rotationally/Magnetically-Driven Winds

The arguments given in the preceding section indicate that physically unreasonable coronal heating rates are required to produce thermally-driven mass loss from T Tauri stars at rates $3 \times 10^{-9} \lesssim \dot{M} \lesssim 10^{-8} M_{\odot} \text{ yr}^{-1}$. However, if T Tauri stars are both rapidly rotating and strongly magnetized, the apparent necessity of a high coronal gas temperature can be circumvented through the inclusion of centrifugal and magnetic forces. To investigate this possibility, consider a steady, axisymmetric outflow in the equatorial plane of a T Tauri star which rotates rigidly with angular frequency $\Omega = \alpha(GM_*/r_0^3)^{1/2}$, where α is a constant. In the absence of magnetic effects, the initial velocity of an isothermal centrifugally-driven wind is approximately

$$V_0 \approx az_c^2 \exp\left\{-\frac{1}{2} - 2z_S\left[1 - \frac{\alpha^2}{2}\left(1 - \frac{1}{z_c^2}\right) - \frac{1}{z_c}\right]\right\} \quad (13)$$

where $a = (kT/\mu)^{1/2}$ is the sound speed,

$$z_c = r_c/r_0 = \frac{1}{2}z_S[1 + (1 - 4\alpha^2/z_S)^{1/2}] \quad (14)$$

is the critical point location, $z_S = (GM_*/2kTr_0)$, and α is restricted to the range $0 \leq \alpha \leq (1 - 1/z_S)^{1/2}$ (Weidelt 1973; Mufson and Liszt 1975; Hartmann and MacGregor 1982b). If a magnetic field is included according to the prescription of Weber and Davis (1967), then

$$V_0 \approx az_c^2 \exp\left\{-\frac{1}{2} - 2z_S\left[1 - \frac{\alpha^2}{2}(z_c^2 - 1) - \frac{1}{z_c}\right]\right\} \quad (15)$$

with

$$z_c \approx 3z_S/(1 + 3\alpha^{2/3}z_S) \quad (16)$$

(Hartmann and MacGregor 1982b). The validity of equations (15) and (16) requires that the stellar magnetic field be strong enough to ensure approximate corotation of the gas throughout the region $1 \leq z \leq z_c$ (cf. Hartmann and MacGregor 1982b).

It is important to note that in order to produce a cool ($T \sim 10^4 \text{ K}$), rotationally and/or magnetically-driven outflow having $\dot{M} \gtrsim 10^{-9} M_{\odot} \text{ yr}^{-1}$ from a T Tauri star, values of $\alpha \sim 1$ are needed. For example, using the stellar parameters given earlier in this section and assuming $r_0 = R_*$, $T = 10^4 \text{ K}$, $N_0 = 10^{12} \text{ cm}^{-3}$, $B_0 = 10 \text{ G}$, it follows from equations (15) and (16) that $z_c \approx 1.07$, $V_0 \approx 0.1a$, and $\dot{M} \approx 1.9 \times 10^{-9} M_{\odot} \text{ yr}^{-1}$ for $\alpha = 0.90$, while for $\alpha = 0.95$, $z_c \approx 1.03$, $V_0 \approx 0.4a$, and $\dot{M} \approx 7.6 \times 10^{-9} M_{\odot} \text{ yr}^{-1}$. The plausibility of the mechanism is therefore dependent upon whether or not T Tauri stars are rapid rotators. A preliminary indication that the answer to this question might be affirmative is provided by the work of Skumanich (1972). He has found that the rotational velocities of solar-type stars decay in time according to $(\text{stellar age})^{-1/2}$; extrapolation of this relation to pre-main-sequence objects would imply rotational velocities $\gtrsim 100 \text{ km s}^{-1}$ for T Tauri stars. Unfortunately, this expectation is not supported by observational results. In particular, Vogel and Kuhi (1981) have measured rotational velocities for 64 pre-main-sequence stars and find that virtually all T Tauri stars with masses $< 1.5 M_{\odot}$ have rotation speeds $\lesssim 25\text{--}35 \text{ km s}^{-1}$. Since this implies $\alpha \sim 0.1\text{--}0.2$ for a T Tauri star with the adopted physical parameters, it is unlikely that winds with $\dot{M} \gtrsim 10^{-9} M_{\odot} \text{ yr}^{-1}$ could be rotationally/magnetically-driven.

iii. Radiatively-Driven Winds

The possibility that mass loss from T Tauri stars is radiatively-driven can be eliminated on the basis of a comparison of the rate of momentum transport by the wind $\dot{M} V_\infty$ with the rate at which momentum can be supplied to the flow by the stellar radiation field L_*/c . If each photon emitted by the star is scattered once in the wind, conservation of momentum indicates that $\dot{M} V_\infty = L_*/c$. For a T Tauri star having $M_* = 0.75 M_\odot$, $R_* = 4 R_\odot$, and $T_{\text{eff}} = 4000$ K, $V_{\text{esc}} = (2GM_*/R_*)^{1/2} \approx 267 \text{ km s}^{-1}$ and $L_* = 1.4 \times 10^{34} \text{ erg s}^{-1} = 3.7 L_\odot$. Assuming $V_\infty = V_{\text{esc}}$ and $3 \times 10^{-9} \lesssim \dot{M} \lesssim 10^{-6} M_\odot \text{ yr}^{-1}$ (DeCampli 1981), the resulting range of wind momentum fluxes is $5.1 \times 10^{24} \lesssim \dot{M} V_\infty \lesssim 1.7 \times 10^{27} \text{ dyne}$, while $L_*/c = 4.7 \times 10^{23} \text{ dyne}$. Since $\dot{M} V_\infty \gg L_*/c$ we conclude that the winds of T Tauri stars cannot be driven by single-scattering radiation pressure alone. It is important to note that this result is a consequence of the low stellar luminosity (L_* is typically $\lesssim 5 L_\odot$ for T Tauri stars), and may be modified if provision is made for the multiple scattering of photons in the flow (Friend and Castor 1983). However, models for the winds of luminous ($L_* \sim 10^5 - 10^6 L_\odot$) early-type stars which include a consistent treatment of this effect are generally characterized by $\dot{M} V_\infty c / L_* \lesssim 10$ (see, e.g., the paper by Friend in this volume), whereas the above analysis suggests that for T Tauri stars $\dot{M} V_\infty c / L_* \gtrsim 10$.

iv. Alfven Wave-Driven Winds

Several authors (DeCampli 1981; Hartmann, Edwards, and Avrett 1982) have suggested that mass loss from T Tauri stars can be driven by the force associated with an outwardly propagating flux of Alfven waves. The applicability of such a mechanism can be examined by using the qualitative wind model described in section II b (iv). For a spherically-symmetric outflow, equation (9) for V_0 together with conservation of mass can be used to derive an approximate expression for the mass loss rate due to a wave-driven wind, (cf. Leer, Holzer, and Flå 1982),

$$\dot{M} \approx 1.825 \times 10^{-13} \left(\frac{r_0}{R_\odot} \right)^{7/2} \left(\frac{M_*}{M_\odot} \right)^{-3/2} \left(\frac{F_{A0}}{10^6 \text{ erg cm}^{-2} \text{ s}^{-1}} \right)^2 \left(\frac{B_0}{\text{gauss}} \right)^{-2} M_\odot \text{ yr}^{-1}, \quad (17)$$

where F_{A0} and B_0 are, respectively, the wave energy flux and magnetic field strength at a reference level $r_0 (= R_*)$. To obtain an estimate of the maximum mass loss rate possible with this mechanism, we first impose the restriction that F_{A0} be less than or equal to the stellar radiative flux $L_*/4\pi R_*^2$. For the hypothetical T Tauri star considered throughout this section, this constraint, implies $F_{A0} \leq 1.45 \times 10^{10} \text{ erg cm}^{-2} \text{ s}^{-1}$. Alternatively, note that for an Alfvénic disturbance with the arbitrarily chosen initial amplitude $\delta B_0 = B_0/\sqrt{10}$, the wave energy flux is simply $F_{A0} \approx (\delta B_0^2/8\pi) \cdot (B_0/\sqrt{4\pi\rho_0}) \approx 1.06 \times 10^9 B_0^3 N_0^{-1/2} \text{ erg cm}^{-2} \text{ s}^{-1}$. Adopting $N_0 = 10^{12} \text{ cm}^{-3}$, it then follows that the requirement $F_{A0} \leq L_*/4\pi R_*^2$ limits the field strength B_0 to values $\lesssim 239 \text{ gauss}$. From equation (17), the corresponding limit on the mass loss rate is $\dot{M} \lesssim 1.32 \times 10^{-7} M_\odot \text{ yr}^{-1}$. This qualitative result is substantiated by detailed model calculations carried out by DeCampli (1981) and Hartmann, Edwards, and Avrett (1982): namely, mass loss rates in the range $\dot{M} \sim 10^{-9} - 10^{-7} M_\odot \text{ yr}^{-1}$ can be obtained for wave energy fluxes F_{A0} of a magnitude such that $4\pi R_*^2 F_{A0} \sim 10^{-1} - 10^0 L_*$. Moreover, although wave dissipation and heating may occur throughout the wind (perhaps giving rise to a variety of optical and ultraviolet emission features; cf. Hartmann *et al.* 1982), the observed terminal velocities appear to be adequately accounted for by models in which the waves propagate without damping (cf. DeCampli 1981). The central (and presently unanswerable) question regarding the applicability of the mechanism to T Tauri star winds is whether or not wave energy fluxes of the required magnitude are physically realizable.

IV. Conclusions

It is apparent from the discussion given in the preceding sections that completely satisfactory theories for mass loss from late-type evolved and pre-main-sequence stars are presently unavailable. It should also be evident that if significant progress is to be made toward the goal of identifying and understanding the physical processes responsible for the winds of cool stars, quantitatively reliable estimates of mass loss rates, terminal velocities, and wind temperatures are necessary. In this regard, it is imperative that the analysis of observational data be carried out using the best available theoretical techniques.

In the case of mass loss from cool giant and supergiant stars, the central problem with which potential driving mechanisms must contend is not just the production of winds having $10^{-11} \lesssim \dot{M} \lesssim 10^{-6} M_{\odot} \text{ yr}^{-1}$. Indeed, for giants and supergiants a substantial portion of the increase in \dot{M} above the solar wind mass loss rate ($\dot{M}_{\odot} \sim 10^{-14} M_{\odot} \text{ yr}^{-1}$) can be attributed to the fact that such stars have larger surface areas; simply scaling the solar wind mass flux density to a star of M supergiant size ($R_{*} \sim 1000 R_{\odot}$) yields $\dot{M} \sim \dot{M}_{\odot} (R_{*}/R_{\odot})^2 \sim 10^{-8} M_{\odot} \text{ yr}^{-1}$. Rather, the most enigmatic property of the observed outflows is that they are generally characterized by expansion velocities which are well below the gravitational escape speed from the stellar surface. The production of a wind having both a large mass loss rate and $V_{\infty} < V_{\text{esc}}$ requires that: (i) most of the energy supplied to the wind by the acceleration mechanism be deposited in the subsonic portion of the flow (to ensure a large \dot{M}); and (ii), the driving force exceed gravity by only a small amount throughout the supersonic portion of the flow (to ensure $V_{\infty} < V_{\text{esc}}$) (cf. Leer and Holzer 1980; Holzer, Flå, and Leer 1983). Virtually all of the mechanisms which have been proposed to account for mass loss from cool, low-gravity stars lack these properties. This difficulty represents perhaps the most formidable obstacle which must be overcome if observation and theory are to be brought into agreement.

An outstanding problem which must be addressed by any theory of mass loss from T Tauri stars is to account for the apparent efficiency with which the driving mechanism operates (cf. DeCampli 1981). Assuming $V_{\infty} \sim V_{\text{esc}}$, the rate at which kinetic energy is transported by the wind is $\frac{1}{2} \dot{M} V_{\infty}^2 \sim GM_{*} \dot{M} / R_{*}$. If the rates of mass loss from T Tauri stars are actually as large as $10^{-8} - 10^{-7} M_{\odot} \text{ yr}^{-1}$, then $\frac{1}{2} \dot{M} V_{\infty}^2 / L_{*} \sim 10^{-2} - 10^{-1}$ for the stellar parameters given at the beginning of section III. The corresponding values of this parameter for the sun, an M supergiant ($\dot{M} \sim 10^{-6} M_{\odot} \text{ yr}^{-1}$, $V_{\infty} \sim 10 \text{ km s}^{-1}$, $L_{*} \sim 10^5 L_{\odot}$), and an O supergiant ($\dot{M} \sim 10^{-6} M_{\odot} \text{ yr}^{-1}$, $V_{\infty} \sim 1000 \text{ km s}^{-1}$, $L_{*} \sim 10^6 L_{\odot}$) are $\sim 10^{-7}$, 10^{-7} , and 10^{-4} , respectively. Because of their intrinsic inefficiency, the mass loss mechanisms considered in section III require excessive energy input rates in order to drive winds with $\dot{M} \gtrsim 10^{-9} M_{\odot} \text{ yr}^{-1}$.

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